

Jet and disk luminosities in tidal disruption events

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12 January 2015

ABSTRACT

Tidal disruption events (TDE) in which a star is devoured by a massive black hole at a galactic center pose a challenge to our understanding of accretion processes. Within a month the accretion rate reaches super-Eddington levels. It then drops gradually over a time scale of a year to sub-Eddington regimes. The initially geometrically thick disk becomes a thin one and eventually an ADAF at very low accretion rates. As such, TDEs explore the whole range of accretion rates and configurations. A challenging question is what the corresponding light curves of these events are. We explore numerically the disk luminosity and the conditions within the inner region of the disk using a fully general relativistic slim disk model. Those conditions determine the magnitude of the magnetic field that engulfs the black hole and this, in turn, determines the Blandford-Znajek jet power. We estimate this power in two different ways and show that they are self-consistent. We find, as expected earlier from analytic arguments (Krolik & Piran 2012), that neither the disk luminosity nor the jet power follows the accretion rate throughout the disruption event. The disk luminosity varies only logarithmically with the accretion rate at super-Eddington luminosities. The jet power follows initially the accretion rate but remains a constant after the transition from super- to sub-Eddington. At lower accretion rates at the end of the MAD phase the disk becomes thin and the jet may stop altogether. These new estimates of the jet power and disk luminosity that do not simply follow the mass fallback rate should be taken into account when searching for TDEs and analysing light curves of TDE candidates. Identification of some of the above mentioned transitions may enable us to estimate better TDE parameters.

Key words: accretion, accretion discs – black hole physics – relativity – methods: numerical – galaxies: jets

1 INTRODUCTION

The surprising X-ray emission from the tidal disruption events (TDE) candidates *Swift* J1644 (Bloom et al. 2011; Levan et al. 2011; Burrows et al. 2011) and *Swift* J2058 (Cenko et al. 2012) led to a re-examination of the accretion processes that take place within these events. 25 years ago, Rees (1988) outlined the basic dynamical processes relevant to TDEs. A star is disrupted by a supermassive black hole. The stellar material is spread out and it returns, after a delay, to the vicinity of the black hole where it forms an accretion disk. If the disrupted stellar material has a uniform distribution of orbital binding energy per unit mass¹, the infall rate onto the accretion disk satisfies: $dM/dt \propto t^{-5/3}$ (Phinney 1989). At the peak accretion rate, the luminosity would be super-Eddington

(Ulmer 1999). Within this disk, gravitational energy is dissipated and the heat is radiated in the usual quasi-thermal fashion. If it is thermally radiated by the accretion disk, the associated temperature would be in the extreme ultra-violet (EUV) or perhaps the soft X-ray band (see e.g. Lodato & Rossi 2011). Consequently, most searches for such events hitherto have been carried out in the EUV region. Indeed several such candidates were found (Gezari et al. 2008; Cappelluti et al. 2009; van Velzen et al. 2011; Arcavi et al. 2014).

In early 2011, Giannios & Metzger (2011) suggested that TDEs would involve ejection of relativistic jet that would give rise to a radio signal when it is slowed down by the surrounding matter. This prediction was verified when the radio emission from *Swift* J1644 was discovered a few months later (Zauderer et al. 2011). However, it turned out that the jet is also a powerful source of X-rays (Bloom et al. 2011; Burrows et al. 2011). These X-rays arise, most likely, from the inner part of the jet, probably due to internal shocks taking place there. A few months later, Cenko et al. (2012) discovered a similar X-ray signature from another TDE candidate

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¹ This assumption that was verified by various numerical simulations, see e.g. Lodato, King & Pringle (2009); Guillochon & Ramirez-Ruiz (2013)

(*Swift* J2058). The latter source also showed an optical emission, most likely the thermal component arising from the accretion disk (Cenko et al. 2012).

The interplay between the (unexpected) nonthermal jet emission and the thermal disk component is intriguing. Over a short period of a few months the accretion rate spans a large range of values, beginning at super-Eddington, turning to sub-Eddington power and then diminishing. Krolik & Piran (2012) addressed this issue using simplified models for both the thermal luminosity and for the jet power. For the former they assumed that the luminosity is capped at the Eddington luminosity at the super-Eddington phase and then it decreases following the mass accretion rate. For the jet power Krolik & Piran (2012) related the Blandford-Znajek (hereafter BZ, Blandford & Znajek 1977; McKinney 2005; Hawley & Krolik 2006) jet power to the pressure at the inner parts of the disk. This pressure is linearly related to the mass accretion rate at the super-Eddington phase. But the pressure is independent of the mass accretion rate during the radiation dominated phase, that follows the super-Eddington phase (Moderski & Sikora 1996). This led Krolik & Piran (2012) to suggest that the jet luminosity will be constant during this phase. At sufficiently low accretion rates the disk becomes gas pressure dominated and the pressure will decrease as the accretion rate decreases. Thus one would expect two transitions in the jet's and disk's light curves. One when the accretion rate drops below Eddington and the other when the disk becomes gas pressure dominated. A third transition would arise at very low accretion rate, once the disk becomes an advection dominated accretion flow (ADAF) (Abramowicz et al. 1988; Narayan & Yi 1995).

Our goal here is to confront these rather simple estimates with a more detailed computation of the accretion disk structure as well as with recent numerical results on the jet luminosity. To this, end we use relativistic slim disk solution (Sądowski 2011) to estimate the accretion disk structure over a large range of accretion rates. The slim disk model generalises the standard thin disks (Novikov & Thorne 1973) towards high accretion rates. It allows for non-Keplerian rotation, radial gradients of pressure, and advection of heat. As a result, when the accretion rate is super-Eddington, the disks are no longer radiatively efficient and their angular momentum profile can significantly depart from Keplerian. The slim disk model uses α viscosity and assumes the accretion rate is constant, i.e., there are no outflows.

Simulations of jet formation from accreting BHs have led to different estimates of the jet luminosity. The early, pioneering simulations of jet formation in general relativistic magnetohydrodynamic (GRMHD) simulations (De Villiers & Hawley 2003; Hawley & Krolik 2006) suggested that the magnetic pressure in the funnel regions of the jet is determined by the thermal pressure of the accretion disk near the BH horizon. It also appeared that the accretion disk structure did not seem to be affected by the presence of the jets or by how strong the jets are (Beckwith et al. 2008). Later, it became clear that both disk thermal pressure and jet magnetic pressure can be much higher than previously simulated in GR and can be as high as the ram pressure of the infalling gas (Tchekhovskoy et al. 2011; Tchekhovskoy & McKinney 2012; McKinney et al. 2012). When the jet pressure reaches this limit, the magnetic field is strong enough to obstruct the accretion of gas onto the BH, and this leads to the formation of a *magnetically arrested disk*, or a MAD (Tchekhovskoy et al. 2011; see Igumenshchev et al. 2003; Igumenshchev 2008 for simulations of MADs in Newtonian relativity and Bisnovatyi-Kogan & Ruzmaikin 1974, 1976; Narayan et al.

2003; Tchekhovskoy 2015 for an analytic consideration). In the following we will combine the simpler jet luminosity estimates based on the disk pressure with the model fits coming from sophisticated numerical simulations.

We begin in §2 examining different methods for estimating the jet power. In §3 we discuss the slim disk model of an accretion flow. We present both the thermal emission and the jet power and we compare the expected behaviour of the jet and disk luminosity for different accretion rates in §4. We examine the implications for these findings to tidal disruption events in §5. We conclude in §6 with a discussion of the limitations of the analysis and possible observational implications.

2 JET POWER

At the order of magnitude level, the Poynting luminosity of a BZ jet that emerges from the vicinity of a black hole can be estimated through a simple dimensional argument. The local magnetic energy density near the black hole is proportional to B^2 , where B is the poloidal field intensity near the horizon. The jet power depends on this energy density and on the area from which the jet emerges. This area is of order $\gtrsim \pi r_g^2$, where r_g is the gravitational radius of the black hole. Combined the overall luminosity can be written as $P_{\text{jet}} = f(a/M) c (B^2/8\pi) (\pi r_g^2)$, where M is the black hole's mass and $a_* \equiv a/M$ its specific angular momentum. The function $f(a_*)$ is dimensionless and we approximate it here as $f(a_*) \approx a_*^2$ (e.g., Tchekhovskoy et al. 2012; Tchekhovskoy 2015). The task now is to better estimate the magnetic field (B) and the area from which the outflow emerges.

2.1 The pressure formula

The jet is in pressure balance with the inner parts of the accretion disk, therefore the strength of the magnetic field is determined by the pressure in the inner disk. Beckwith et al. (2008) have demonstrated that the magnetic pressure near the horizon is generally bounded above by the maximal pressure in the equatorial plane near the inner edge of the disk, p_{max} , and it is bounded below by the magnetic pressure at that location. Thus, we can estimate the strength of the magnetic energy density using p_{max} , the maximal pressure near the edge of the disk:

$$B^2/8\pi = \beta_B p_{\text{max}}, \quad (1)$$

where the factor β_B is an unknown dimensionless factor of order unity. A second factor is the size of the region from which the jet emerges. We use the radius of the innermost stable circular orbit (ISCO), R_{ISCO} , to characterise the size of the inner disk, which is comparable but smaller than the position of the maximal pressure. Together, we can write the BZ luminosity as:

$$P_{\text{jet,P}} \equiv \beta_B \pi R_{\text{ISCO}}^2 p_{\text{max}} a_*^2 c \lesssim \pi R_{\text{ISCO}}^2 p_{\text{max}} a_*^2 c. \quad (2)$$

The unknown coefficient β_B will be determined in Section 2.3 by comparing with the jet power obtained in numerical simulations.

2.2 The H/R formula

As in the previous section, we will balance the magnetic pressure pressure in the jet funnel against the pressure of the accretion disk. The vertical force balance approximately gives,

$$\frac{p}{\rho} \approx \Omega_K^2 H^2 = \frac{V_K^2 H^2}{R^2}, \quad (3)$$

where p and ρ , are the midplane pressure and the corresponding density of the disk, respectively, Ω_K and V_K are the Keplerian velocities at radius R , and H is disk half-thickness. Applying the vertical equilibrium at the radius of the pressure maximum, R , to eq. 2, we get,

$$P_{\text{jet,H/R}} \equiv \beta'_B \pi V_K^2 \Sigma H a_*^2 c, \quad (4)$$

where we introduced the surface density $\Sigma = 2\rho H$, and allow the coefficient β'_B to differ from β_B . The rest mass conservation requires,

$$\dot{M} = 2\pi R \Sigma V_r, \quad (5)$$

where \dot{M} is the accretion rate and V_r is the absolute value of the radial velocity. Using this formula we get,

$$P_{\text{jet,H/R}} = \frac{\beta'_B}{2} \dot{M} c^2 \frac{V_K^2}{V_r c} \frac{H}{R} a_*^2. \quad (6)$$

For radiatively inefficient accretion flows, the magnitude of the radial velocity V_r is comparable to the Keplerian velocity V_K and does not depend on the accretion rate. We may therefore approximate in this case eq. 6 and write,

$$P_{\text{jet,H/R}} \propto \dot{M} c^2 \frac{H}{R} a_*^2. \quad (7)$$

2.3 Normalisation

Numerical simulations of jets in both optically thin and thick, radiatively inefficient MAD disks (Tchekhovskoy et al. 2011; Tchekhovskoy & McKinney 2012; McKinney et al. 2012, 2013) provide the missing scaling factor and give (Tchekhovskoy 2015),

$$P_{\text{jet,H/R,RIAF}} = 1.3 \dot{M} c^2 \frac{H/R}{0.3} a_*^2. \quad (8)$$

Because of the taken assumptions, this formula for the jet power is valid only for radiatively inefficient disks — ADAFs or super-critical (or super-Eddington) disks, i.e., in the limit of lowest and highest accretion rates. To estimate the power of the jet in disks with not so extreme accretion rates, one has to use more general formulation (eq. 2 or eq. 6).

All the three formulae are expected to give the same estimate of the jet power. To satisfy this condition, we choose the coefficients β_B (eq. 2) and β'_B (eq. 6) so that the jet power estimates agree with eq. 8 for $\dot{M} \gg \dot{M}_{\text{Edd}}$, and finally get,

$$P_{\text{jet,P}} = 4.0 \pi R_{\text{ISCO}}^2 p_{\text{max}} a_*^2 c, \quad (9)$$

and

$$P_{\text{jet,H/R}} = 0.16 \dot{M} c^2 \frac{V_K^2}{V_r c} \frac{H/R}{0.3} a_*^2. \quad (10)$$

3 DISK MODEL

To model an accretion disk we use the general relativistic slim disk solutions of Sadowski (2011). The slim disk model generalises the standard thin disk (Shakura & Sunyaev 1973; Novikov & Thorne 1973) to arbitrary accretion rates. It allows for non-Keplerian rotation and advective cooling by photons trapped in the flow. It assumes constant accretion rate, i.e., it does not allow for outflows, and it adopts the α prescription for viscosity. The slim disks reduce in the limit of small accretion rates to the standard thin, relativistic, Keplerian disk.

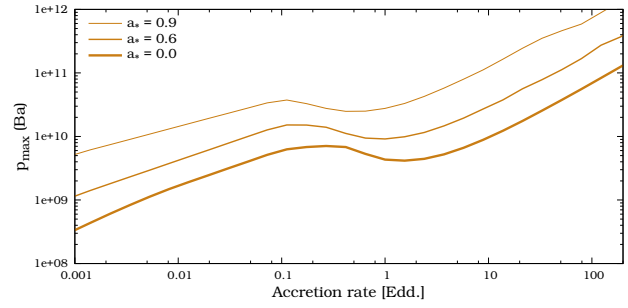


Figure 1. The maximal pressure in the disk as a function of the accretion rate for BH spins $a_* = 0.3, 0.6$ and 0.9 and $M = 10^7 M_\odot$.

The system has a characteristic luminosity, the Eddington luminosity, L_{Edd} :

$$L_{\text{Edd}} = \frac{4\pi G M c}{\kappa_{\text{es}}} = 1.25 \times 10^{38} \frac{M}{M_\odot} \text{ erg/s}. \quad (11)$$

This corresponds to the Eddington accretion rate, here defined as:

$$\dot{M}_{\text{Edd}} = \frac{1}{\eta} \frac{L_{\text{Edd}}}{c^2} = \frac{1}{\eta} \frac{4\pi G M}{c \kappa_{\text{es}}} = 2.44 \times 10^{18} \frac{M}{M_\odot} \text{ g/s}, \quad (12)$$

where we put the efficiency of a thin disk around a non-rotating BH, $\eta = 0.057$, and $\kappa_{\text{es}} = 0.4 \text{ cm}^2/\text{g}$. Once the accretion rate is near and above \dot{M}_{Edd} , photons do not have enough time to diffuse out of the disk and a fraction of them is advected with the flow. This extra advective cooling modifies the structure of the disk. In particular, the disk radiates less efficiently and its luminosity scales as (Paczynski 1980),

$$L \approx L_{\text{Edd}} \left(1 + \log \frac{\dot{M}}{\dot{M}_{\text{Edd}}} \right). \quad (13)$$

The power of the jet depends on the parameters of the underlying accretion flow. The pressure formula for the jet power (eq. 9) is parametrised in terms of the maximal total pressure in the equatorial plane, p_{max} . Fig. 1 presents p_{max} as a function of the accretion rate for three values of BH spin: $a_* = 0.0$ (thickest), 0.6 , and 0.9 (thinnest line), assuming a BH mass $M = 10^7 M_\odot$. For the lowest accretion rates, $\dot{M} \lesssim 0.1 \dot{M}_{\text{Edd}}$, the disk is gas pressure dominated and the pressure at fixed radius is expected to follow $\dot{M}^{4/5}$ (Shakura & Sunyaev 1973). However, the radius of the pressure maximum is not fixed and the profiles of this quantity follow this dependence only qualitatively. When accretion rate exceeds $\sim 0.1 \dot{M}_{\text{Edd}}$ the disk is radiation pressure dominated. The standard thin disk theory predicts that the pressure at the equatorial plane at a given radius is independent of the accretion rate. This explains the flattening of p_{max} profiles around $0.1 \dot{M}_{\text{Edd}}$. Because the pressure maximum is not at a fixed radius and because the advective cooling gradually kicks in when approaching \dot{M}_{Edd} , the p_{max} slightly varies with \dot{M} . For super-critical (exceeding \dot{M}_{Edd}) accretion rates, the disk enters the slim disk regime and the maximal pressure is proportional to the accretion rate.

The H/R -based estimates of the jet power (eqs. 8 and 10) depend on the disk thickness which we parametrise by the maximal disk opening angle, $\Theta_H = \arctan(H/R)$. Fig. 2 presents profiles of disk thickness for various accretion rates, BH spins and masses. The top panel corresponds to a non-rotating BH. For $\dot{m} = \dot{M}/\dot{M}_{\text{Edd}} = 0.3$ the thickest radiation pressure dominated region is located between $R = 10$ and 100 . Disk thickness is determined there by the local radiative flux. Because both the flux and the vertical component of the gravity are proportional to the BH mass, the

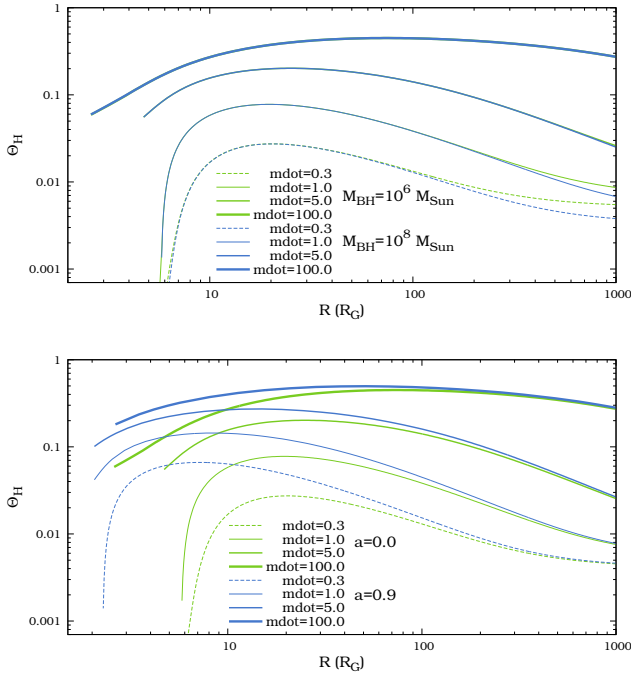


Figure 2. Radial profiles of disk thickness (defined as the opening angle $\Theta_H = \text{atan}(H/R)$) for various accretion rates. The top panel compares two BH masses (10^6 vs $10^8 M_\odot$) assuming zero BH spin, while the bottom one compares two values of BH spin ($a_* = 0$ vs 0.9) for a BH mass $10^7 M_\odot$.

disk opening angle does not depend on the BH mass. Further out, where gas pressure dominates, the BH mass has a slight impact on disk thickness. The extent of the radiation pressure dominated region increases with accretion rate and reaches $R \gtrsim 1000$ for highly super-critical accretion rates.

The bottom panel of Fig. 2 depicts the disk opening angle for a fixed BH mass $M = 10^7 M_\odot$ and two values of BH spin, $a_* = 0.0$ (green) and 0.9 (blue lines). The disks (and hence the thickness profiles) extend more inward for the rotating BH. This reflects the fact that the radius of the innermost stable circular orbit decreases with increasing BH spin. At the same time, for a fixed accretion rate, disks around rotating BH have a higher maximal thickness than their non-spinning counterparts. This reflects the higher accretion efficiency that results in a higher luminosity.

Finally, in Fig. 3 we plot the maximal disk opening angle as a function of accretion rate for three BH spins. When the disk is gas pressure dominated at all radii ($\dot{M} \lesssim 0.001 \dot{M}_{\text{Edd}}$) the maximal disk thickness is located far from the BH and it does not depend on its spin — the lines therefore coincide. In the intermediate, radiation pressure dominated, regime the disk thickness increases with the BH spin, as discussed above. For the radiatively inefficient (slim) regime, the maximal disk thickness saturates at $\Theta_H \approx 0.3$. This reflects the fact that the thickness of advection dominated accretion flows depends only on the ratio of the radiation to gas pressure, and approaches the limiting value $\Theta_H \approx 0.3$ for radiation pressure dominated super-critical disks (Narayan & Yi 1994; Vieira et al. 2015).

4 JET AND DISK LUMINOSITIES

The standard model of thin disks (Shakura & Sunyaev 1973) predicts that energy liberated is proportional to the mass accretion, \dot{M} ,

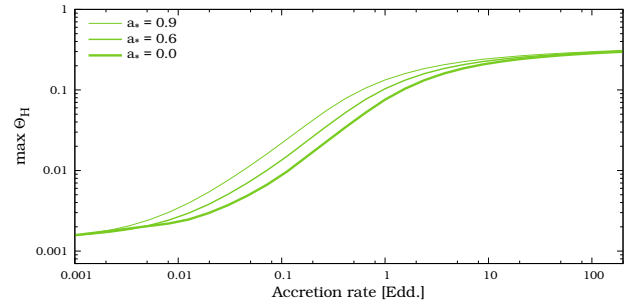


Figure 3. The maximal opening angle, Θ_H , as a function of the accretion rate for BH spins $a_* = 0.3, 0.6$ and 0.9 .

and determined by the accretion efficiency, η ,

$$L = \eta \dot{M} c^2. \quad (14)$$

For non-rotating BH $\eta = 0.057$. As discussed above, once the accretion rate approaches and exceeds the Eddington limit, the efficiency is decreased. These facts are reflected in the disk luminosity profile shown in Fig. 4, corresponding to a BH spin $a_* = 0.6$ and a BH mass $M_{\text{BH}} = 10^7 M_\odot$, with the red line. As long as $\dot{M} < \dot{M}_{\text{Edd}}$, the disk luminosity increases linearly with the accretion rate. Once this limit has been exceeded, the luminosity grows roughly with logarithm of \dot{M} (eq. 13).

We now apply the jet power formulae derived in Section 2 and estimate the jet power for each accretion rate using the corresponding slim disk solution as the underlying disk model.

The pressure formula (eq. 9) reflects the fact that the disk and jet pressures balance each other. Knowing the disk thermal pressure we may therefore estimate the jet magnetic pressure, and therefore the magnetic flux in the jet. This quantity, together with the known BH spin, provides an estimate of the jet power. Its dependence on mass accretion rate is shown in Fig. 4 with the blue line. For a fixed BH mass and spin eq. 9 depends only on the maximal value of the disk thermal pressure, and therefore the profile of the jet power estimated this way resembles the profile of the corresponding maximal disk pressure (Fig. 1). The jet power in the super-Eddington regime grows proportionally to the accretion rate, remains roughly constant in the radiation-pressure dominated thin disk regime ($0.1 < \dot{M}/\dot{M}_{\text{Edd}} < 1$), and follows $\dot{M}^{4/5}$ for gas-pressure dominated thin disks. These properties result in very powerful jet power estimates (exceeding the disk luminosity by 2-3 orders of magnitude) for the thinnest disks. This is inconsistent with observations (see, e.g., Fender et al. 2004; Russell et al. 2011) and suggests that the assumptions behind the pressure formula break down in this regime.

The other formula for the jet power (the H/R formula, eq. 10), similarly uses the disk thermal pressure as the proxy for the jet magnetic pressure and magnetic flux at the horizon. However, the disk pressure is replaced with the disk thickness using the vertical equilibrium equation, and the radial velocity V_r and the disk thickness H are introduced. These are taken directly from the numerical solutions of slim disks. The corresponding jet power is plotted with the green line in Fig. 4. It coincides with the pressure formula for $\dot{M} > \dot{M}_{\text{Edd}}$ and stays close for lower accretion rates. This fact proves that both formulae properly identify the disk thermal pressure, although the H/R formula does it indirectly. From now on we will not distinguish between these two ways of estimating the jet power and use eq. 9 as the proxy.

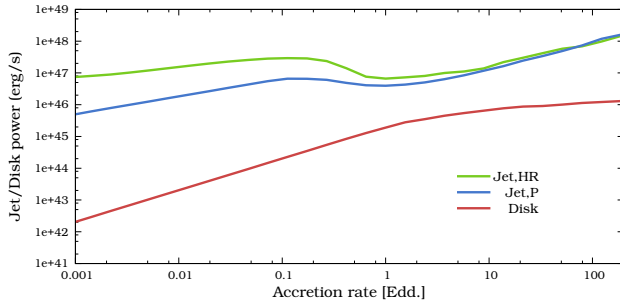


Figure 4. Jet (green and blue) and disk (red line) luminosities for $M_{\text{BH}} = 10^7 M_{\odot}$ and BH spin $a_* = 0.6$ a function of accretion rate. The green line corresponds to eq. 10. The blue line shows the jet power calculated according to eq. 9, which we take below as our fundamental jet power estimate.

In Fig. 5 we plot the jet power and its efficiency,

$$\eta = \frac{L_{\text{jet}}}{\dot{M}c^2},$$

as a function of accretion rate for three values of BH spin with solid blue and dotted black lines, respectively. The profiles of the jet power for all spins agrees with the discussion above. The jet power is expected to grow roughly with BH spin squared, therefore, the higher the spin, the larger is the jet power. In the limit of highest accretion rates, the efficiency calculated according to both formulae for spin $a_* = 0.9$ is close to $\eta = 1$, what reflects the fact that the efficiencies of jets observed in GRMHD simulations may be very high, and even exceed 1 for the highest spins and magnetically arrested disks (Tchekhovskoy et al. 2011; Tchekhovskoy & McKinney 2012).

The jet efficiency remains roughly constant as long as we are in the radiatively inefficient regime ($\dot{M} > \dot{M}_{\text{Edd}}$). Below that limit the efficiency increases with decreasing accretion rate. Because the disk thermal pressure stays roughly constant in the radiation-pressure dominated, thin disk regime, when the accretion rate changes by an order of magnitude, or more, the efficiency must increase by that amount. Furthermore, the pressure dependence on accretion rate for gas-pressure dominated disks is $P \propto \dot{M}^{4/5}$, what further increases the efficiency. Such extreme efficiencies for the thinnest disks ($\eta \approx 100$ for $0.001\dot{M}_{\text{Edd}}$) are unphysical and inconsistent with the lack of observed jets in high/soft state of Galactic X-ray binaries.

Which assumptions we took break down in the sub-Eddington regime? Why are the jets not there? It is probably the assumption that β_B that determines the ratio of the magnetic pressure to the disk pressure remains constant and is independent of the dramatic changes that take place in the disk when it transits from a thick to thin disk and from radiation pressure to gas pressure dominated regime. In particular, one expects the radial disk velocity to decrease and the large-scale magnetic flux to diffuse outward faster than it is being advected inward by the accretion flow (Lubow et al. 1994; Guilet & Ogilvie 2012, 2013; see however Rothstein & Lovelace 2008). In addition, one expects for thin disks a larger gap between the inner part of the disk and the black hole. As the accretion rate decreases, less and less material is within this gap and the thin disk may not be able to drag the necessary magnetic field to the vicinity of the black hole.

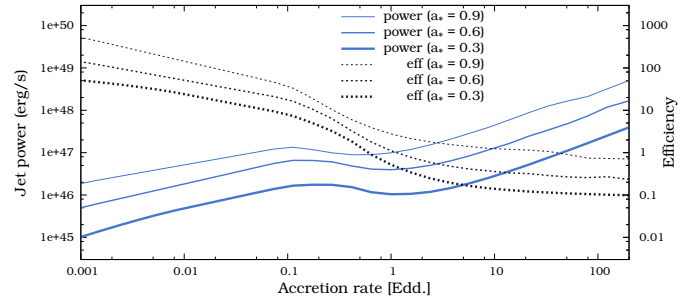


Figure 5. Jet power (blue lines, left y-axis) and its efficiency (given by eq. (4) and shown with dotted black lines, right y-axis) for $M_{\text{BH}} = 10^7 M_{\odot}$ and BH spins $a_* = 0.3, 0.6$ and 0.9 .

5 JET AND DISK IN TIDAL DISRUPTION EVENTS.

We turn now to apply the previous result to a TDE. We consider a star with mass M_* (measured in solar masses) that is disrupted by a massive black hole (BH). We approximate the main sequence mass-radius relation by $R_* \approx R_{\odot} M_*^{(1-\xi)}$; $\xi \approx 0.2$ for $0.1 < M_* \leq 1$, but increases to ≈ 0.4 for $1 < M_* < 10$ (Kippenhahn & Weigert 1994). Finally we define k as the apsidal motion constant (determined by the star's radial density profile) and f is its binding energy in units of GM_*^2/R_* (Phinney 1989). In numerical estimates, we scale k/f to the value for fully radiative stars, 0.02, because this is a reasonable approximation for main sequence stars with $0.4M_{\odot} < M < 10M_{\odot}$ (Kippenhahn & Weigert 1994).

Assuming that the mass accretion rate follows the fall back time of stellar material onto the central black hole, and that the disrupted star has a uniform distribution in orbital binding energy per unit mass, matter returns to the region near the pericenter radius at a rate $\dot{M} \propto (t/t_0)^{-5/3}$ (Phinney 1989). The characteristic timescale t_0 for initiation of this power-law accretion rate is the orbital period for the most bound matter (Lodato, King & Pringle 2009; Krolik & Piran 2012):

$$t_0 \approx 15 \times 10^5 M_*^{(1-3\xi)/2} M_{\text{BH},7}^{1/2} \left(\frac{k/f}{0.02} \right)^{1/2} \text{ s}. \quad (15)$$

Using this time scale we calibrate the maximal accretion rate² as:

$$\dot{M}_{\text{peak}} \sim 0.3 \times 10^{27} M_*^{(1+3\xi)/2} M_{\text{BH},7}^{-1/2} (k/f)^{-1/2} \text{ g s}^{-1}. \quad (16)$$

We convolve now the previous estimates of jet luminosity and disk power with the accretion rate evolution to obtain the expected light curves. These are shown for $10^6 M_{\odot}$ and $10^7 M_{\odot}$ BHs with $a_* = 0.6$ spin in Fig. 6. A quick inspection reveals that, as expected, even with the more detailed calculations, the thermal disk light curves (red lines) follow more or less the simple estimates of Krolik & Piran (2012). The disk luminosity is roughly constant at short time scales (when $\dot{M} > \dot{M}_{\text{Edd}}$). After the transition to $\dot{M} < \dot{M}_{\text{Edd}}$, which takes place roughly at $t/t_0 = 6$ and 20 for $10^7 M_{\odot}$ and $10^6 M_{\odot}$ BHs, respectively, the disk luminosity decreases proportionally to the mass accretion rate. For a higher BH mass (dotted lines) the qualitative behaviour remains the same. However, because the initial accretion rate (eq. 16) is lower, the disk enters earlier the sub-Eddington and thin disk phases.

² See however Shiokawa et al. (2015) for caution concerning the onset of accretion in TDEs and the possibility of a lower maximal accretion rate that takes place at a later moment.

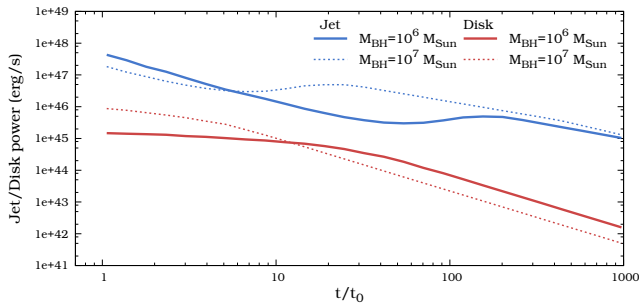


Figure 6. Disk luminosity (red) and jet power (blue line) as a function of time for a tidal disruption event of a $M_* = 1$ star by a $M_{\text{BH}} = 10^6 M_\odot$ (solid) or $10^7 M_\odot$ (dotted lines) BH. Other parameters: $M_* = 1$, $k/f = 0.02$, and $\xi = 0.2$.

The jet luminosity evolution is also consistent with predictions of Krolik & Piran (2012). Initially, as long as the disk is super-Eddington, the jet power (denoted with blue lines) decreases proportionally to the accretion rate. Once it enters the sub-critical, radiatively efficient regime, the jet power hardly changes reflecting the fact that the disk pressure is not sensitive to the accretion rate. This changes once the disk becomes gas pressure dominated. Our simplistic formulae predict further decrease of the jet power, although not as steep as in the super-critical regime. However, as discussed earlier, thin disks are not likely to sustain strong magnetic flux at the BH and therefore should produce less efficient jets than we predict.

6 DISCUSSION AND CONCLUSIONS

We have used numerical modelling to estimate the jet power and the accretion disk luminosity as a function of an accretion rate over a very wide range of accretion rates. Our accretion disk model use the relativistic slim disk solutions of Sądowski (2011). It switches from a thick disk for accretion rates above Eddington to a thin disk below. We have used two different expressions for the jet power, one based on the maximal pressure and another, motivated by numerical simulation, based on the H/R ratio. Remarkably, both formulae give comparable results.

Our results confirm earlier expectations (Krolik & Piran 2012) that were based on simple scaling arguments. According to these expectations the jet power and the disk luminosity do not simply follow the accretion rate. Specifically, for super-Eddington accretion the jet power follows the accretion rate but the disk bolometric luminosity varies only logarithmically with it. At sub-Eddington accretion the jet power becomes roughly constant while the disk bolometric luminosity follows the accretion rate. A puzzling behaviour in our solution is the appearance of a very large efficiency (jet power vs. accretion rate) at low accretion luminosity. While puzzling we note that this is not impossible in principle as the BZ jets gets its power from the rotational energy of the central black hole and as such could, in principle at least, operate even with lower accretion rate (provided that it is surrounded by strong enough magnetic fields). However, as we discuss shortly below, it is likely that the jet power switches off completely in this low accretion rate regime.

It is interesting to note that in our model the jet power is higher by at least one order or magnitude than the disk luminosity. This holds as long as a jet exists for the whole range of accretion rates

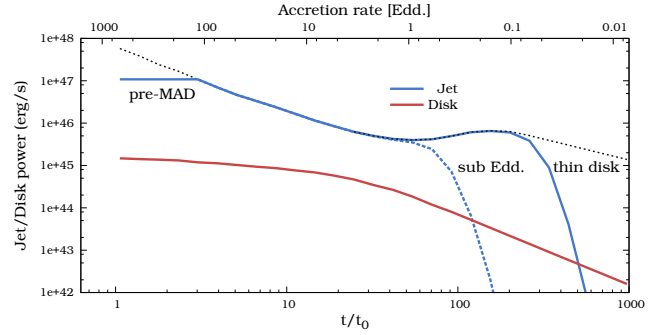


Figure 7. Similar to Fig. 6 but only for $M_{\text{BH}} = 10^6 M_\odot$. The labels above the top axis denote the accretion rate in Eddington units at given time. The plateau before $t = 3t_0$ schematically reflects the pre-MAD stage when the accretion rate is too large to form magnetically arrested disk. The jet power was arbitrarily damped for $t > 200t_0$ to reflect the fact that the disk enters the thin disk regime ($\dot{M} < 0.1\dot{M}_{\text{Edd}}$) and is not expected to produce efficient jet at all. The dotted blue line reflects corresponding decrease in the jet power if the transition takes place already at \dot{M}_{Edd} .

and time scales considered. This may not be related directly to the question whether it is easier or more difficult to detect the jet emission or the disk emission as those depend also on the spectral energy distribution of these two components and on sensitivity and coverage of different detectors in the relevant spectral regimes.

It should be stressed that three possibly important effects have not been taken into consideration in our work. First, we assumed (implicitly) that the magnetic fields needed for the BZ mechanism to operate are there right from the beginning. This might not be the case in TDEs in which the infalling material is not highly magnetised initially.³ Thus it may take some time to build this magnetic field and in this case we expect a constant (or even increasing) jet power at early times (Tchekhovskoy et al. 2014). This effect was marked in Fig. 7 as a solid horizontal line marked “pre-MAD”. In other transient astrophysical source, core-collapse gamma-ray bursts, this stage lasts for the most duration of the event: in fact, then the MAD onset plausibly marks the end of the burst (Tchekhovskoy & Giannios 2015).

A second potential effect that we have ignored here is the possibility of the jet switching off when the disk makes a transition from thick to thin (Zauderer et al. 2013; Tchekhovskoy et al. 2014). This transition takes place shortly after the accretion rate makes the transition from super to sub-Eddington. As mentioned earlier such a turn off of the jet can arise from the fact that the weaker accretion rate does not push the magnetic field sufficiently close to the black hole, or that the large-scale magnetic flux diffuses outward more effectively than it is dragged inward by the disk. This possible effect was marked in Fig. 7 by arbitrary damping the jet power for $\dot{M} < 0.1\dot{M}_{\text{Edd}}$ what indicates that the jet power may not satisfy the simple curve in this region. Note that in the future, as mass accretion rate drops even further, $\dot{M} \lesssim 0.01\dot{M}_{\text{Edd}}$, the accretion disk is expected to transition to a geometrically-thick radiatively-inefficient ADAF, which can cause re-launching of the jets and X-ray emission from Sw 1644 (Tchekhovskoy et al. 2014).

Finally, we assumed that all of the mass fallback ends up

³ The energetics of Sw J1644 requires the presence of a large-scale magnetic flux that exceeds by 3–5 orders of magnitude the magnetic flux through a main-sequence star. Such a large magnetic flux could be captured by the debris stream from a pre-existing accretion disk (Tchekhovskoy et al. 2014), as recently seen in numerical simulations (Kelley et al. 2014).

reaching the black hole. The processes of gas circularisation are not well-understood, and it is possible that the fraction of gas that reaches the hole depends on the Eddington ratio, and this can cause additional deviations from simple power-law scalings in time. Note also that the scaling vs time of intensity in a particular detector might deviate from a power-law due to the shifting spectrum of the source coming in and out of the detector bandpass, which might further complicate the structure of the detected lightcurves.

To conclude, we remark on a few potential observational implications. First, we note that for all reasonable values of parameters and for all accretion rates the jet power (as long as the jet exists) is at least one order of magnitude larger than the disk luminosity (see, e.g., Fig. 4). This is consistent with observations that jet power exceeds accretion disk luminosity in blazars and that the accretion flow in these systems is in the magnetically-arrested disk regime (Rawlings & Saunders 1991; Zamaninasab et al. 2014; Ghisellini et al. 2014; Zdziarski et al. 2015). Moreover, if the jet is relativistic its radiation would be beamed and enhanced further. Secondly, we note again that both the jet power (and its corresponding X-ray emission) and the disk luminosity do not follow in a simple manner the mass accretion rate. This should be taken into account when searching observationally for TDEs or when analysing the observed light curves of TDE candidates. Particularly interesting is the possibility (Krolik & Piran 2012) that some of the above mentioned transitions and in particular the transition from super- to sub-Eddington accretion or the formation of a thin disk could be identified. This would provide significant new independent information on the parameters of the TDE and in particular on the masses of massive black hole.

7 ACKNOWLEDGEMENTS

We thank Marek Abramowicz for initiating this collaboration. T.P. was supported by an ERC advanced grant (GRBs) and by the I-CORE Program of the Planning and Budgeting Committee and The Israel Science Foundation (grant No 1829/12). A.S. and A.T. acknowledge support for this work by NASA through Einstein Postdoctoral Fellowships PF4-150126 and PF3-140115, respectively, awarded by the Chandra X-ray Center, which is operated by the Smithsonian Astrophysical Observatory for NASA under contract NAS8-03060.

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